

Soft Terms from Strings *

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ABSTRACT

We study the structure of the soft SUSY-breaking terms obtained from 4-D Strings under the assumption of dilaton/moduli dominance in the process of SUSY breaking. We first analyze in detail the dilaton-dominated limit because of its finiteness properties and phenomenological predictivity, and second, we consider the new features appearing when several moduli fields contribute to SUSY breaking. In particular, we discuss in detail the case of symmetric Abelian orbifolds. Although some qualitative features indeed change in the multimoduli case with respect to the dilaton dominance one, the most natural mass relations at low-energy, $m_l < m_q \simeq M_g$, are still similar. Only in some very specific limits these relations might be reversed. We also study the presence of tachyons pointing out that their possible existence may be, in some cases, an interesting advantage in order to break extra gauge symmetries. Finally, we compute explicitly the μ and B parameters in the context of the mechanism for generating a “ μ -term” by the Kähler potential, as naturally implemented in orbifolds. It leads to the prediction $|tg\beta| = 1$ at the String scale, independently of the Goldstino direction. It is worth noticing that in this scheme the dilaton-dominated case, where there is no free parameters, is excluded since it is not consistent with the measured value of the top-quark mass. In this connection, low-energy charge and color breaking minima are also discussed.

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1 INTRODUCTION AND SUMMARY

In the last four years there has been much activity in trying to obtain information about the structure of soft Super-symmetry (SUSY)-breaking terms in effective $N = 1$ theories coming from four-dimensional (4-D) Strings [1]. The basic idea is to identify some $N = 1$ chiral fields whose auxiliary components could break SUSY by acquiring a vacuum expectation value (VEV). Natural candidates in 4-D Strings are 1) the complex dilaton field $S = \frac{1}{g^2} + ia$ which is present in any 4-D String and 2) the moduli fields T^i, U^i which parametrize the size and shape of the compactified variety in models obtained by compactification of a ten-dimensional heterotic String. It is not totally unreasonable to think that some of these fields may play an important role in SUSY breaking. To start with, if String models are to make any sense, these fields should be strongly affected by non-perturbative phenomena. They are massless in perturbation theory and non-perturbative effects should give them a mass to avoid deviations from the equivalence principle and other phenomenological problems. Secondly, these fields are generically present in large classes of 4-D models (the dilaton in all of them). Finally, the couplings of these fields to charged matter are suppressed by powers of the Planck mass, which makes them natural candidates to constitute the SUSY-breaking “hidden sector” which is assumed to be present in phenomenological models of low-energy SUSY.

The important point in this assumption of locating the seed of SUSY-breaking in the dilaton/moduli sectors, is that it leads to some interesting relationships among different soft terms [2-7] which could perhaps be experimentally tested. This general analysis was applied in particular to the gaugino condensation scenario in ref.[2], whereas in refs.[3-7] no special assumption was made about the possible origin of SUSY breaking.

In ref.[5] a systematic discussion of the structure of soft terms which may be obtained under the assumption of dilaton/moduli-dominated SUSY breaking in some classes of 4-D Strings was presented, with particular emphasis on the case of Abelian $(0, 2)$ orbifold models [8]. It was mostly considered a situation in which only the dilaton S and an “overall modulus T ” field contribute to SUSY breaking. In fact, actual 4-D Strings like orbifolds contain several T_i and U_i moduli. Generic $(0, 2)$ orbifold models contain three T_i moduli fields (only Z_3 has 9 and Z_4, Z'_6 have 5) and a maximum of three (“complex structure”) U_i fields. The use of an overall modulus T is equivalent to the assumption that the three T_i fields of generic orbifold models contribute exactly the same to SUSY breaking and the rest do not contribute. In the absence of further dynamical information it is reasonable to expect similar contributions from *all* the moduli although not necessarily exactly the same. Thus it is natural to ask what changes if one relaxes the overall modulus hypothesis and works with the multimoduli case [7]. This is one of the purposes of the present talk.

The second one is to analyze in more detail the dilaton-dominated limit, where only the dilaton field contributes to SUSY breaking [4, 5]. This is a very interesting possibility not only due to phenomenological reasons, as universality of the soft terms, but also to theoretical arguments. In this connection it has recently been realized [9, 10, 11] that the boundary conditions $-A = M_{1/2} = \sqrt{3}m$ of dilaton dominance coincide with some boundary conditions considered by Jones, Mezzinescu and Yau in 1984 [12] in a complete different context. They found that those same boundary conditions maintain the (two-loop) finiteness properties of certain $N = 1$ SUSY theories [13]. This coincidence is in principle quite surprising since we did not bother about the loop corrections when extracting these boundary conditions from the dilaton-dominance assumption. Also, effective $N = 1$ field theories from Strings do not in general fulfil the finiteness requirements. It has also been noticed [10] that this coincidence could be related to an underlying $N = 4$ structure of the dilaton Lagrangian and that the dilaton-dominated boundary conditions could appear as a fixed point of renormalization group equations [10, 14]. This could perhaps be an indication that at least some of the possible soft terms obtained in the present scheme could have a more general relevance, not necessarily linked to a particular form of the tree-level Lagrangian.

In section 2 we present an analysis of the effects of several moduli on the results obtained for soft terms. In the multimoduli case several parameters are needed to specify the Goldstino direction in the dilaton/moduli space, in contrast with the overall modulus case where the relevant information is contained in just one angular parameter θ . The presence of more free parameters leads to some loss of predictivity for the soft terms. This predictivity is recovered and increased in the case of dilaton dominance where the soft terms, eq.(8), are independent of the 4-D String considered and fulfil the low-energy mass relations given by eq.(9). Also we show that, even in the multimoduli case, in some schemes there are certain sum-rules among soft terms, eq.(13), which hold independently of the Goldstino direction. The presence of these sum rules cause that, *on average* the *qualitative* results in the dilaton-dominated case still apply. Specifically, if one insists e.g. in obtaining scalar masses heavier than gauginos (something not possible in the dilaton-dominated scenario), this is possible in the multimoduli case, but the sum-rules often force some of the scalars to get negative squared mass. If we want to avoid this, we have to stick to gaugino masses bigger than (or of order) the scalar masses. This would lead us back to the qualitative results obtained in dilaton dominance. Let us notice however that in very specific limits, which will be discussed below, these results might be modified. In the case of standard model 4-D Strings the tachyonic behaviour may be particularly problematic, since charge and/or colour could be broken. In the case of GUTs constructed from Strings, it may just be the signal of GUT symmetry breaking.

However, even in this case one expects the same order of magnitude results for observable scalar and gaugino masses and hence the most natural mass relations *at low-energy* are still similar to the dilaton dominance ones.

Another topic of interest is the B parameter, the soft mass term which is associated to a SUSY-mass term $\mu H_1 H_2$ for the pair of Higgses $H_{1,2}$ in the Minimal Supersymmetric Standard Model (MSSM). Compared to the other soft terms, the result for the B parameter is more model dependent. Indeed, it depends not only on the dilaton/moduli-dominance assumption but also on the particular mechanism which could generate the associated “ μ term” [1]. An interesting possibility to generate such a term is the one suggested in refs.[15, 16] in which it was pointed out that in the presence of certain bilinear terms in the Kähler potential an effective μ term of order the gravitino mass, $m_{3/2}$, is naturally generated. Interestingly enough, such bilinear terms in the Kähler potential do appear in String models and particularly in Abelian orbifolds. In section 3 we compute the μ and B parameters³ as well as the soft scalar masses of the charged fields which could play the role of Higgs particles in such Abelian orbifold schemes. We find the interesting result that, *independently of the Goldstino direction* in the dilaton/moduli space, one gets the prediction $|tg\beta| = 1$ at the String scale. On the other hand, if we consider the interesting Goldstino direction where only the dilaton breaks SUSY, the whole soft terms and the μ parameter depend only on the gravitino mass. Imposing the phenomenological requirement of correct electroweak breaking we arrive to the remarkable result that the whole SUSY spectrum is completely determined with no free parameters. Unfortunately, this direction is not consistent with the measured value of the top-quark mass. In this connection, an interesting comment about low-energy charge and color breaking minima in the dilaton-dominated limit can be found at the end of the section.

A few comments before closing up this summary are in order. First of all we are assuming here that the seed of SUSY breaking propagates through the auxiliary fields of the dilaton S and the moduli T_i, U_i fields. However attractive this possibility might be, it is fair to say that there is no compelling reason why indeed no other fields in the theory could participate. Nevertheless the present scheme has a certain predictivity due to the relative universality of the couplings of the dilaton and moduli. Indeed, the dilaton has universal and model-independent couplings which are there independently of the 4-D String considered. The moduli T_i, U_i fields are less universal, their number and structure depend on the type of compactification considered. However, there are thousands of different (0,2) models with different particle content which share the same T_i, U_i moduli structure. For example, the moduli structure of a given Z_N orbifold is the same for all the thousands of (0,2) models one can construct from it by doing different embeddings and adding discrete Wilson lines. So, in this sense, although not really universal, there are large classes of models with identical T_i, U_i couplings. This is not the case of generic charged matter fields whose number and couplings are completely out of control, each individual model being in general completely different from any other. Thus assuming dilaton/moduli dominance in the SUSY-breaking process has at least the advantage of leading to specific predictions for large classes of models whereas if charged matter fields play an important role in SUSY breaking we will be forced to a model by model analysis, something which looks out of reach.

Another point to remark is that we will use the tree level forms for both the gauge kinetic function and the Kähler potential. One-loop corrections to these functions have been computed in refs.[17] and [18] respectively in some classes of 4-D Strings (orbifold models) and their effects on the soft terms have also been studied [5, 19, 20] and could be included in the analysis without difficulty. In fact, the effects of these one-loop corrections will in general be negligible except for those corners of the Goldstino directions in which the tree-level soft terms vanish. However, as we will see below, this situation would be a sort of fine-tuning. More worrisome are the possible non-perturbative String corrections to the Kähler and gauge kinetic functions. We have made use in our orbifold models of the known tree-level results for those functions. If the non-perturbative String corrections turn out to be important, it would be impossible to make any prediction about soft terms unless we know all the relevant non-perturbative String dynamics, something which looks rather remote (although perhaps not so remote as it looked one year ago!).

2 SOFT TERMS

2.1 General structure: the multimoduli case

We are going to consider $N = 1$ SUSY 4-D Strings with m moduli $T_i, i = 1, \dots, m$. Such notation refers to both T -type and U -type (Kähler class and complex structure in the Calabi-Yau language) fields. In addition there will be charged matter fields C_α and the complex dilaton field S . In general we will be considering (0,2) compactifications and thus the charged fields do not need to correspond to 27s of E_6 .

Before further specifying the class of theories that we are going to consider a comment about the total number of moduli is in order. We are used to think of large numbers of T and U -like moduli due to the fact that in (2,2) (E_6) compactifications there is a one to one correspondence between moduli and charged fields. However, in the case of (0,2) models with arbitrary gauge group (which is the case of phenomenological interest) the number of moduli

³The results for B corresponding to the possibility of generating a small μ term from the superpotential [16] can also be found, for the multimoduli case under consideration, in ref.[7]. They are more model dependent.

is drastically reduced. For example, in the standard $(2, 2)$ Z_3 orbifold there are 36 moduli T_i , 9 associated to the untwisted sector and 27 to the fixed points of the orbifold. In the thousands of $(0, 2)$ Z_3 orbifolds one can construct by adding different gauge backgrounds or doing different gauge embeddings, only the 9 untwisted moduli remain in the spectrum. The same applies to models with U -fields. This is also the case for compactifications using $(2, 2)$ minimal superconformal models. Here all singlets associated to twisted sectors are projected out when proceeding to $(0, 2)$ [21]. So, as these examples show, in the case of $(0, 2)$ compactifications the number of moduli is drastically reduced to a few fields. In the case of generic Abelian orbifolds one is in fact left with only three T-type moduli T_i ($i = 1, 2, 3$), the only exceptions being Z_3 , Z_4 and Z'_6 , where such number is 9, 5 and 5 respectively. The number of U -type fields in these $(0, 2)$ orbifolds oscillates between 0 and 3, depending on the specific example. Specifically, $(0, 2)$ $Z_2 \times Z_2$ orbifolds have 3 U fields, the orbifolds of type $Z_4, Z_6, Z_8, Z_2 \times Z_4, Z_2 \times Z_6$ and Z'_{12} have just one U field and the rest have no untwisted U -fields. Thus, apart from the three exceptions mentioned above, this class of models has at most 6 moduli, three of T -type (always present) and at most three of U -type. In the case of models obtained from Calabi-Yau type of compactifications a similar effect is expected and only one T -field associated to the overall modulus is guaranteed to exist in $(0, 2)$ models.

We will consider effective $N = 1$ supergravity (SUGRA) Kähler potentials of the type:

$$K(S, S^*, T_i, T_i^*, C_\alpha, C_\alpha^*) = -\log(S + S^*) + \hat{K}(T_i, T_i^*) + \tilde{K}_{\bar{\alpha}\beta}(T_i, T_i^*) C^{*\bar{\alpha}} C^\beta + (Z_{\alpha\beta}(T_i, T_i^*) C^\alpha C^\beta + h.c.). \quad (1)$$

The first piece is the usual term corresponding to the complex dilaton S which is present for any compactification whereas the second is the Kähler potential of the moduli fields, where we recall that we are denoting the T - and U -type moduli collectively by T_i . The greek indices label the matter fields and their kinetic term functions are given by $\tilde{K}_{\bar{\alpha}\beta}$ and $Z_{\alpha\beta}$ to lowest order in the matter fields. The last piece is often forbidden by gauge invariance in specific models although it may be relevant in some cases as discussed in section 3. The complete $N = 1$ SUGRA Lagrangian is determined by the Kähler potential $K(\phi_M, \phi_M^*)$, the superpotential $W(\phi_M)$ and the gauge kinetic functions $f_a(\phi_M)$, where ϕ_M generically denotes the chiral fields S, T_i, C_α . As is well known, K and W appear in the Lagrangian only in the combination $G = K + \log |W|^2$. In particular, the (F-part of the) scalar potential is given by

$$V(\phi_M, \phi_M^*) = e^G \left(G_M K^{M\bar{N}} G_{\bar{N}} - 3 \right), \quad (2)$$

where $G_M \equiv \partial_M G \equiv \partial G / \partial \phi_M$ and $K^{M\bar{N}}$ is the inverse of the Kähler metric $K_{\bar{N}M} \equiv \partial_{\bar{N}} \partial_M K$.

The crucial assumption now is to locate the origin of SUSY breaking in the dilaton/moduli sector. Then, plugging eq.(1) into eq.(2), the bosonic soft SUSY-breaking terms can be computed. Applying the standard SUGRA formulae [22] to the most general case where the moduli and matter metrics are not diagonal we obtain:

$$m_{\bar{\alpha}\beta}^2 = m_{3/2}^2 \tilde{K}_{\bar{\alpha}\beta} - \bar{F}^{\bar{i}} (\partial_{\bar{i}} \partial_{\bar{j}} \tilde{K}_{\bar{\alpha}\beta} - \partial_{\bar{i}} \tilde{K}_{\bar{\alpha}\gamma} \tilde{K}^{\gamma\bar{\delta}} \partial_{\bar{j}} \tilde{K}_{\bar{\delta}\beta}) F^j, \quad (3)$$

$$A'_{\alpha\beta\gamma} = F^S K_S h_{\alpha\beta\gamma} + F^i \left[\hat{K}_i h_{\alpha\beta\gamma} + \partial_i h_{\alpha\beta\gamma} - \left(\tilde{K}^{\delta\bar{\rho}} \partial_i \tilde{K}_{\bar{\rho}\alpha} h_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right], \quad (4)$$

where $m_{\bar{\alpha}\beta}^2$ and $A'_{\alpha\beta\gamma}$ are the soft mass matrix and the soft trilinear parameters respectively (corresponding to unnormalized charged fields), $h_{\alpha\beta\gamma}$ is a (un-rescaled) renormalizable Yukawa coupling involving three charged chiral fields and $F^S = e^{G/2} K_{S\bar{S}}^{-1} G_{\bar{S}}$, $F^i = e^{G/2} \hat{K}^{i\bar{j}} G_{\bar{j}}$ are the dilaton and moduli auxiliary fields. Notice that, after normalizing the fields to get canonical kinetic terms, the first piece in eq.(3) will lead to universal diagonal soft masses but the second piece will generically induce off-diagonal contributions. Concerning the A -parameters, notice that we have not factored out the Yukawa couplings as usual, since proportionality is not guaranteed. Indeed, although the first term in $A'_{\alpha\beta\gamma}$ is always proportional in flavour space to the corresponding Yukawa coupling, the same thing is not necessarily true for the other terms. In this section we are going to consider the case of diagonal metric both for the moduli and the matter fields⁴. Then $\hat{K}(T_i, T_i^*)$ will be a sum of contributions (one for each T_i), whereas $\tilde{K}_{\bar{\alpha}\beta}$ will be taken of the diagonal form $\tilde{K}_{\bar{\alpha}\beta} \equiv \delta_{\bar{\alpha}\beta} \tilde{K}_\alpha$.

Let us take the following parametrization for the VEV's of the dilaton and moduli auxiliary fields

$$\begin{aligned} G_{S\bar{S}}^{1/2} F^S &= \sqrt{3} m_{3/2} \sin \theta e^{-i\gamma_S}, \\ G_{i\bar{i}}^{1/2} F^i &= \sqrt{3} m_{3/2} \cos \theta e^{-i\gamma_i} \Theta_i, \end{aligned} \quad (5)$$

where $\sum_i \Theta_i^2 = 1$ and $e^G = m_{3/2}^2$ is the gravitino mass-squared. The angle θ and the Θ_i just parametrize the direction of the goldstino in the S, T_i field space. We have also allowed for the possibility of some complex phases

⁴An extensive analysis of the off-diagonal case in specific orbifold constructions, including the calculation of the soft terms and their effects on flavour changing neutral currents (FCNC), can be found in ref.[7].

γ_S, γ_i which could be relevant for the CP structure of the theory. This parametrization has the virtue that when we plug it in the general form of the SUGRA scalar potential eq.(2), its VEV (the cosmological constant) vanishes by construction. Notice that such a phenomenological approach allows us to ‘reabsorb’ (or circumvent) our ignorance about the (nonperturbative) S - and T_i - dependent part of the superpotential, which is responsible for SUSY breaking.

It is now a straightforward exercise to compute the bosonic soft SUSY-breaking terms in this class of theories. Plugging eq.(5) into eqs.(3,4) one finds the following results (we recall that we are considering here a diagonal metric for the matter fields):

$$m_\alpha^2 = m_{3/2}^2 \left[1 - 3 \cos^2 \theta (\hat{K}_{\bar{i}i})^{-1/2} \Theta_i e^{i\gamma_i} (\log \tilde{K}_\alpha)_{\bar{i}j} (\hat{K}_{\bar{j}j})^{-1/2} \Theta_j e^{-i\gamma_j} \right] ,$$

$$A_{\alpha\beta\gamma} = -\sqrt{3} m_{3/2} \left[e^{-i\gamma_S} \sin \theta - e^{-i\gamma_i} \cos \theta \Theta_i (\hat{K}_{\bar{i}i})^{-1/2} \left(\hat{K}_i - \sum_{\delta=\alpha,\beta,\gamma} (\log \tilde{K}_\delta)_i + (\log h_{\alpha\beta\gamma})_i \right) \right] . \quad (6)$$

The above scalar masses and trilinear scalar couplings (where we have factorized out the Yukawa coupling as usual) correspond to charged fields which have already been canonically normalized.

Physical gaugino masses M_a for the canonically normalized gaugino fields are given in general by $M_a = F^M [\log(Ref_a)]_M$. Since the tree-level gauge kinetic function is given for any 4-D String by $f_a = k_a S$, where k_a is the Kac-Moody level of the gauge factor, the result for tree-level gaugino masses is independent of the moduli sector and is simply given by:

$$M \equiv M_a = m_{3/2} \sqrt{3} \sin \theta e^{-i\gamma_S} . \quad (7)$$

As we mentioned above, the parametrization of the auxiliary field VEV’s was chosen in such a way to guarantee the automatic vanishing of the VEV of the scalar potential ($V_0 = 0$). If the value of V_0 is not assumed to be zero the above formulae (5-7) are modified in the following simple way. One just has to replace $m_{3/2} \rightarrow C m_{3/2}$, where $|C|^2 = 1 + V_0/3m_{3/2}^2$. In addition, the formula for m_α^2 gets an additional contribution given by $2m_{3/2}^2(|C|^2 - 1) = 2V_0/3$.

The soft term formulae above (6, 7) are in general valid for any compactification as long as we are considering diagonal metrics. In addition one is tacitly assuming that the tree-level Kähler potential and f_a -functions constitute a good approximation. The Kähler potentials for the moduli are in general complicated functions. Before going into specific classes of Superstring models, it is worth studying the interesting limit $\cos \theta = 0$, corresponding to the case where the dilaton sector is the source of all the SUSY breaking (see eq.(5)).

2.2 The $\cos \theta = 0$ (dilaton-dominated) limit

Since the dilaton couples in an universal manner to all particles, *this limit is quite model independent*. Using eqs.(6,7) one finds the following simple expressions for the soft terms which are independent of the 4-D String considered

$$m_\alpha = m_{3/2} ,$$

$$M_a = \pm \sqrt{3} m_{3/2} ,$$

$$A_{\alpha\beta\gamma} = -M_a , \quad (8)$$

where, from the limits on the electric dipole moment of the neutron, we have imposed $\gamma_S = 0 \bmod \pi$.

This dilaton-dominated scenario [4, 5] is attractive for its simplicity and for the natural explanation that it offers to the universality of the soft terms. Actually, universality is a desirable property not only to reduce the number of independent parameters in the MSSM, but also for phenomenological reasons, particularly to avoid FCNC.

Because of the simplicity of this scenario, the low-energy predictions are quite precise [23, 5, 24]. Since scalars are lighter than gauginos at the String scale, at low-energy ($\sim M_Z$), gluino, slepton and (first and second generation) squark mass relations turn out to be⁵

$$M_g : m_Q : m_u : m_d : m_L : M_e \simeq 1 : 0.94 : 0.92 : 0.92 : 0.32 : 0.24 . \quad (9)$$

Although squarks and sleptons have the same soft mass, at low-energy the former are much heavier than the latter because of the gluino contribution to the renormalization of their masses.

2.3 Orbifold compactifications

To illustrate some general features of the multimoduli case we will concentrate here on the case of generic (0,2) symmetric Abelian orbifolds. As we mentioned above, this class of models contains three T -type moduli and (at most)

⁵The phenomenology of SUSY breaking by the dilaton in the context of a flipped $SU(5)$ model was also studied in ref.[25].

three U -type moduli. We will denote them collectively by T_i , where e.g. $T_i = U_{i-3}$; $i = 4, 5, 6$. For this class of models the Kähler potential has the form [26]

$$K(\phi, \phi^*) = -\log(S + S^*) - \sum_i \log(T_i + T_i^*) + \sum_\alpha |C_\alpha|^2 \Pi_i(T_i + T_i^*)^{n_\alpha^i}. \quad (10)$$

Here n_α^i are fractional numbers usually called “modular weights” of the matter fields C_α . For each given Abelian orbifold, independently of the gauge group or particle content, the possible values of the modular weights are very restricted. For a classification of modular weights for all Abelian orbifolds see ref.[3]. As a matter of fact, the Kähler potentials which appear in the large- T limit of Calabi-Yau compactifications [27] and 4-D fermionic Strings [28] are quite close to the above one. Thus the results that we will obtain below will probably be more general than just for orbifold compactifications.

Using the particular form (10) of the Kähler potential and eqs.(6,7) we obtain the following results⁶ for the scalar masses, gaugino masses and soft trilinear couplings:

$$\begin{aligned} m_\alpha^2 &= m_{3/2}^2 (1 + 3 \cos^2 \theta \, n_\alpha^i \cdot \vec{\Theta}^2), \\ M &= \sqrt{3} m_{3/2} \sin \theta e^{-i\gamma_S}, \\ A_{\alpha\beta\gamma} &= -\sqrt{3} m_{3/2} (\sin \theta e^{-i\gamma_S} + \cos \theta \sum_{i=1}^6 e^{-i\gamma_i} \Theta^i \omega_{\alpha\beta\gamma}^i), \end{aligned} \quad (11)$$

where we have defined :

$$\omega_{\alpha\beta\gamma}^i = (1 + n_\alpha^i + n_\beta^i + n_\gamma^i - Y_{\alpha\beta\gamma}^i) ; Y_{\alpha\beta\gamma}^i = \frac{h_{\alpha\beta\gamma}^i}{h_{\alpha\beta\gamma}} 2ReT_i. \quad (12)$$

Notice that neither the scalar nor the gaugino masses have any explicit dependence on S or T_i , they only depend on the gravitino mass and the goldstino angles. This is one of the advantages of a parametrization in terms of such angles. Although in the case of the A -parameter an explicit T_i -dependence may appear in the term proportional to $Y_{\alpha\beta\gamma}^i$, it disappears in several interesting cases [7]. With the above information we can now analyze the structure of soft terms available for Abelian orbifolds.

1) Universality of soft terms

In the dilaton-dominated case ($\cos \theta = 0$) the whole soft terms are universal. However, in general, they show a lack of universality due to the modular weight dependence (see eqs.(11,12)).

2) Soft masses

In the multimoduli case, depending on the goldstino direction, tachyons may appear. For $\cos^2 \theta \geq 1/3$, one has to be very careful with the goldstino direction if one is interested in avoiding tachyons. Nevertheless, as we will discuss below, having a tachyonic sector is not necessarily a problem, it may even be an advantage, so one should not disregard this possibility at this point.

Consider now three particles $C_\alpha, C_\beta, C_\gamma$ coupling through a Yukawa $h_{\alpha\beta\gamma}$. They may belong both to the untwisted (**U**) sector or to a twisted (**T**) sector, i.e. couplings of the type **UUU**, **UTT**, **TTT**. Then, using the above formulae, one finds [7] that in general for *any choice* of goldstino direction

$$m_\alpha^2 + m_\beta^2 + m_\gamma^2 \leq |M|^2 = 3m_{3/2}^2 \sin^2 \theta. \quad (13)$$

Notice that if we insist in having a vanishing gaugino mass, the sum-rule (13) forces the scalars to be either all massless or at least one of them tachyonic. Nevertheless we should not forget that tachyons, as we already mentioned above, are not necessarily a problem, but may just show us an instability.

3) Gaugino versus scalar masses

In the multimoduli case on average the scalars are lighter than gauginos but there may be scalars with mass bigger than gauginos. Eq.(13) tells us that this can only be true at the cost of having some of the other three scalars with *negative* squared mass. This may have diverse phenomenological implications depending what is the particle content of the model, as we now explain in some detail:

3-a) Gaugino versus scalar masses in standard model 4-D Strings

Let us suppose we insist in e.g., having tree-level gaugino masses lighter than the scalar masses. If we are dealing with a String model with gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y \times G$ this is potentially a disaster. Some observable particles, like Higgses, squarks or sleptons would be forced to acquire large VEV's (of order the String scale). For example, the scalars associated through the Yukawa coupling $H_2 Q_L u_L^c$, which generates the mass of the u -quark, must

⁶This analysis was also carried out for the particular case of the three diagonal moduli T_i in ref.[19] and [20] in order to obtain unification of gauge coupling constants and to analyze FCNC constraints respectively. Some particular multimoduli examples were also considered in ref.[6].

fulfil the above sum-rule (13). If we allow e.g. the scalars H_2 , Q_L to be heavier than gauginos, then u_L^c will become tachyonic breaking charge and color. However, tachyons may be helpful if the particular Yukawa coupling does not involve observable particles. They could break extra gauge symmetries and generate large masses for extra particles. We recall that standard-like models in Strings usually have too many extra particles and many extra U(1) interactions. Although the Fayet-Iliopoulos mechanism helps to cure the problem [29], the existence of tachyons is a complementary solution.

We thus see that, for standard model Strings, if we want to avoid charge and colour-breaking minima (or VEV's of order the String scale for the Higgses⁷), we should grosso modo come back to a situation with gauginos heavier than scalars. Thus the low-energy phenomenological predictions of the multimoduli case are similar to those of the dilaton-dominated scenario (see subsect.2.2): due to the sum-rule the tree-level observable scalars are always lighter than gauginos

$$m_\alpha < M . \quad (14)$$

Now, at low-energy ($\sim M_Z$), gluino, slepton and (first and second generation) squark mass relations turn out to be

$$m_l < m_q \simeq M_g , \quad (15)$$

where gluinos are slightly heavier than scalars. This result is qualitatively similar to the dilaton dominance one, in spite of the different set of (non-universal) soft scalar masses, because the low-energy scalar masses are mainly determined by the gaugino loop contributions. The only exception are the sleptons masses, which do not feel the important gluino contribution, and therefore can get some deviation from the result of eq.(9).

As emphasized in [5] there is however a way to get scalars heavier than gauginos, even in the overall modulus case, if all the observable particles have overall modular weight $n_\alpha = -1$ and $\sin \theta \rightarrow 0$ (i.e. in the moduli-dominated limit). Then, at tree-level, $M \rightarrow 0$ and $m_\alpha \rightarrow 0$ if the different moduli participate in the SUSY breaking in almost *exactly* the same way, i.e. the overall modulus situation. Including String loop corrections to K and f_a can yield scalars heavier than gauginos [5]

$$m_\alpha > M_a \quad (16)$$

and the low-energy spectrum can be reversed with respect to the above one (in the case of $\sin \theta$ sufficiently small as to produce $m_\alpha \gg M_a$)

$$M_g < m_l \simeq m_q . \quad (17)$$

The physical masses of squarks and sleptons are almost degenerate because the universality of soft scalar masses at high-energy is not destroyed by the gluino contribution to the mass renormalization, which is now very small. Notice however that this possibility of obtaining scalars heavier than gauginos is a sort of fine-tuning. In the absence of a more fundamental theory which tells us in what direction the goldstino angles point, one would naively say that the most natural possibility would be to assume that all moduli contribute to SUSY breaking in more or less (but *not* exactly⁸) the same amount.

We just saw how, in the context of standard model Strings, the results for soft terms are qualitatively similar to the dilaton dominance ones if we want to avoid the breaking of charge and colour conservation. There is however a loophole in the above analysis. Up to now we have assumed that the masses of the observable fermions arise through renormalizable Yukawa couplings. If we give up that assumption and allow the existence of non-renormalizable Yukawa couplings generating masses for the observable particles (e.g. $H_2 Q_L u_L^c < \phi \dots \phi >$), then new sum-rules would apply to the full set of fields in the coupling and the above three-particle sum-rules could be violated. In particular, observable scalars would be allowed to be heavier than gauginos, possibly at the price of having some tachyon among the (standard model singlet) ϕ fields. Then qualitative results different from the ones of the dilaton dominance case may be obtained

In this respect, it is easy to find explicit examples of orbifold sectors yielding scalar masses bigger than gaugino masses even at the tree-level. From eq.(11) we see that always $m_\alpha < m_{3/2}$ and therefore scalars heavier than gauginos can be obtained if the constraint

$$\cos^2 \theta > 2/3 \quad (18)$$

is fulfilled. Let us consider e.g. the case of the Z_8 orbifold with an observable particle in the twisted sector \mathbf{T}_{θ^6} . The modular weight associated to that sector is $n_{\theta^6} = (1/4, 3/4, 0, 0)$ and therefore (see eq.(11))

$$m_{\theta^6}^2 = m_{3/2}^2 \left[1 - 3 \cos^2 \theta \left(\frac{1}{4} \Theta_1^2 + \frac{3}{4} \Theta_2^2 \right) \right] . \quad (19)$$

⁷For a possible way-out to this problem, allowing the possibility of scalars heavier than gauginos, see ref.[30].

⁸For an explicit example of this, using gaugino condensation, see ref.[31].

For the particular values $\cos^2 \theta = 5/6$, $\Theta_1 = \Theta_2 = 0$ one gets $m_{\theta^6}^2 = m_{3/2}^2$, $M^2 = m_{3/2}^2/2$.

In spite of the new possibilities offered by the multimoduli extension, one typically finds that, unless very particular choices for the goldstino angles are chosen, the masses of scalar and gauginos are still of the same order and therefore at low-energy eq.(15) is typically still valid, the only difference being that now squarks will be slightly heavier than gluinos. To reverse the situation (i.e. eq.(17)) we would need $m_\alpha \gg M_a$. This can be obtained in the limit $\sin \theta \rightarrow 0$, i.e. $M \rightarrow 0$. However, there may be a phenomenological problem in this case. Experimental bounds on gluino mass imply $M > 50$ GeV which only can be obtained for a large $m_{3/2}$ but this would yield a large $m_\alpha \sim m_{3/2}$. In general one must be careful to avoid m_α bigger than 1 TeV, spoiling the solution to the gauge hierarchy problem.

3-b) Gaugino versus scalar masses in GUT 4-D Strings

What it turned out to be a potential disaster in the case of standard model Strings may be an interesting advantage in the case of String-GUTs. In this case it could well be that the negative squared mass may just induce gauge symmetry breaking by forcing a VEV for a particular scalar (GUT-Higgs field) in the model. The latter possibility provides us with interesting phenomenological consequences. Here the breaking of SUSY would directly induce further gauge symmetry breaking. An explicit example of this situation can be found in ref.[7].

In summary, the situation concerning gaugino versus scalar masses is as follows. If any of the physical quark-lepton Yukawas come from renormalizable terms, the sum rules leads us to a distribution of soft terms in such a way that gaugino masses are generically bigger than those of scalars (otherwise charge and/or colour would be broken). For a possible exception see footnote 5. If the physical Yukawas come all from non-renormalizable terms the constraints coming from the sum rules may be avoided, possibly allowing standard model singlets to become tachyonic. However, even in this case one expects the same order of magnitude results for scalar and gaugino masses and hence the most natural (slepton-squark-gluino) mass relations *at low-energy* will be similar to the ones of the dilaton-dominated case eq.(15) as showed in point 3-a. Only in the particular limit of very small $\sin \theta$ this situation might be reversed.

3 THE B PARAMETER AND THE μ PROBLEM

It was pointed out in refs.[15, 16] that terms in a Kähler potential like the one proportional to $Z_{\alpha\beta}$ in eq.(1) can naturally induce a μ -term for the C_α fields of order $m_{3/2}$ after SUSY breaking, thus providing a rationale for the size of μ . From eqs.(1,2) and from the fermionic part of the SUGRA lagrangian one can check that a SUSY mass term $\mu_{\alpha\beta} C_\alpha C_\beta$ and a scalar term $B_{\alpha\beta}(C_\alpha C_\beta) + h.c.$ are induced upon SUSY breaking in the effective low-energy theory (here the kinetic terms for $C_{\alpha,\beta}$ have not still been canonically normalized)

$$\mu_{\alpha\beta} = m_{3/2} Z_{\alpha\beta} - \bar{F}^{\bar{i}} \partial_{\bar{i}} Z_{\alpha\beta} , \quad (20)$$

$$B_{\alpha\beta} = 2m_{3/2}^2 Z_{\alpha\beta} + m_{3/2} F^i \left[\partial_i Z_{\alpha\beta} - \left(\tilde{K}^{\delta\bar{\rho}} \partial_i \tilde{K}_{\bar{\rho}\alpha} Z_{\delta\beta} + (\alpha \leftrightarrow \beta) \right) \right] - m_{3/2} \bar{F}^{\bar{i}} \partial_{\bar{i}} Z_{\alpha\beta} \\ - \bar{F}^{\bar{i}} F^j \left[\partial_j \partial_{\bar{i}} Z_{\alpha\beta} - \left(\tilde{K}^{\delta\bar{\rho}} \partial_j \tilde{K}_{\bar{\rho}\alpha} \partial_{\bar{i}} Z_{\delta\beta} + (\alpha \leftrightarrow \beta) \right) \right] . \quad (21)$$

Notice that, as in the case of the A -terms and the corresponding Yukawa couplings (see subsection 2.1), $B_{\alpha\beta}$ is not necessarily proportional to $\mu_{\alpha\beta}$.

Recently it has been suggested that terms of the type $Z_{\alpha\beta} C_\alpha C_\beta + h.c.$ may appear in the Kähler potential of some Calabi-Yau type compactifications [4]. It has also been explicitly shown [32] that they appear in orbifold models. Let us consider the case in which e.g., due to gauge invariance, there is only one possible μ -term (and correspondingly one B term) associated to a pair of matter fields C_1, C_2 . This is e.g. the case of the MSSM. If we introduce the abbreviations

$$L^Z \equiv \log Z , \quad L^\alpha \equiv \log \tilde{K}_\alpha , \quad X \equiv 1 - \sqrt{3} \cos \theta e^{i\gamma_i} \Theta_i (\hat{K}_{\bar{i}i})^{-1/2} L_i^Z , \quad (22)$$

using eqs.(20,21) the μ and B parameters are given by

$$\mu = m_{3/2} (\tilde{K}_1 \tilde{K}_2)^{-1/2} Z X , \quad (23)$$

$$B = m_{3/2} X^{-1} \left[2 + \sqrt{3} \cos \theta (\hat{K}_{\bar{i}i})^{-1/2} \Theta_i \left(e^{-i\gamma_i} (L_i^Z - L_i^1 - L_i^2) - e^{i\gamma_i} L_i^Z \right) \right. \\ \left. + 3 \cos^2 \theta (\hat{K}_{\bar{i}i})^{-1/2} \Theta_i e^{i\gamma_i} \left(L_i^Z (L_j^1 + L_j^2) - L_i^Z L_j^Z - L_{ij}^Z \right) (\hat{K}_{\bar{j}j})^{-1/2} \Theta_j e^{-i\gamma_j} \right] , \quad (24)$$

where we are assuming that the moduli on which $\tilde{K}_1(T_i, T_i^*)$, $\tilde{K}_2(T_i, T_i^*)$ and $Z(T_i, T_i^*)$ depend have diagonal metric, which is the relevant case we are going to discuss. The above μ and B (where we have factorized out the μ term as usual) parameters correspond now to charged fields which have already been canonically normalized.

If the value of V_0 is not assumed to be zero, one just has to replace $\cos \theta \rightarrow C \cos \theta$ in eqs.(22,23,24), where C is given below eq.(7). In addition, the formula for B gets an additional contribution given by $m_{3/2} X^{-1} 3(C^2 - 1)$.

As mentioned above, it has recently been shown that the untwisted sector of orbifolds with at least one complex-structure field U possesses the required structure $Z(T_i, T_i^*)C_1C_2 + h.c.$ in their Kähler potentials [32]. Specifically, the Z_N orbifolds based on Z_4, Z_6, Z_8, Z'_{12} and the $Z_N \times Z_M$ orbifolds based on $Z_2 \times Z_4$ and $Z_2 \times Z_6$ do all have a U -type field in (say) the third complex plane. In addition the $Z_2 \times Z_2$ orbifold has U fields in the three complex planes. In all these models the piece of the Kähler potential involving the moduli and the untwisted matter fields $C_{1,2}$ in the third complex plane has the form

$$K(T_i, T_i^*, C_1, C_2) = K'(T_l, T_l^*) - \log((T_3 + T_3^*)(U_3 + U_3^*) - (C_1 + C_1^*)(C_2^* + C_2)) \quad (25)$$

$$\simeq K'(T_l, T_l^*) - \log(T_3 + T_3^*) - \log(U_3 + U_3^*) + \frac{(C_1 + C_2^*)(C_1^* + C_2)}{(T_3 + T_3^*)(U_3 + U_3^*)} . \quad (26)$$

The first term $K'(T_l, T_l^*)$ determines the (not necessarily diagonal) metric of the moduli $T_l \neq T_3, U_3$ associated to the first and second complex planes. The last term describes an $SO(2, n)/SO(2) \times SO(n)$ Kähler manifold ($n = 4$ if we focus on just one component of C_1 and C_2) parametrized by T_3, U_3, C_1, C_2 . If the expansion shown in (26) is performed, on one hand one recovers the well known factorization $SO(2, 2)/SO(2) \times SO(2) \simeq (SU(1, 1)/U(1))^2$ for the submanifold spanned by T_3 and U_3 (which have therefore diagonal metric to lowest order in the matter fields), whereas on the other hand one can easily identify the functions $Z, \tilde{K}_1, \tilde{K}_2$ associated to C_1 and C_2 :

$$Z = \tilde{K}_1 = \tilde{K}_2 = \frac{1}{(T_3 + T_3^*)(U_3 + U_3^*)} . \quad (27)$$

Plugging back these expressions in eqs.(23,24,22) one can compute μ and B for this interesting class of models [7]:

$$\mu = m_{3/2} (1 + \sqrt{3} \cos \theta (e^{i\gamma_3} \Theta_3 + e^{i\gamma_6} \Theta_6)) , \quad (28)$$

$$B\mu = 2m_{3/2}^2 (1 + \sqrt{3} \cos \theta (\cos \gamma_3 \Theta_3 + \cos \gamma_6 \Theta_6) + 3 \cos^2 \theta \cos(\gamma_3 - \gamma_6) \Theta_3 \Theta_6) . \quad (29)$$

In addition, we recall from eq.(11) that the soft masses are

$$m_{C_1}^2 = m_{C_2}^2 = m_{3/2}^2 (1 - 3 \cos^2 \theta (\Theta_3^2 + \Theta_6^2)) . \quad (30)$$

In general, the dimension-two scalar potential for $C_{1,2}$ after SUSY breaking has the form

$$V_2(C_1, C_2) = (m_{C_1}^2 + |\mu|^2)|C_1|^2 + (m_{C_2}^2 + |\mu|^2)|C_2|^2 + (B\mu C_1 C_2 + h.c.) . \quad (31)$$

In the specific case under consideration, from eqs.(28,29,30) we find the remarkable result that the three coefficients in $V_2(C_1, C_2)$ are equal, i.e.

$$m_{C_1}^2 + |\mu|^2 = m_{C_2}^2 + |\mu|^2 = B\mu . \quad (32)$$

so that $V_2(C_1, C_2)$ has the simple form

$$V_2(C_1, C_2) = B\mu (C_1 + C_2^*)(C_1^* + C_2) . \quad (33)$$

Although the common value of the three coefficients in eq.(32) depends on the Goldstino direction via the parameters $\cos \theta, \Theta_3, \Theta_6, \dots$ (see expression of $B\mu$ in eq.(29)), we stress that the equality itself and the form of V_2 hold *independently of the Goldstino direction*. The only constraint that one may want to impose is that the coefficient $B\mu$ be non-negative, which would select a region of parameter space. For instance, if one neglects phases, such requirement can be written simply as

$$(1 + \sqrt{3} \cos \theta \Theta_3)(1 + \sqrt{3} \cos \theta \Theta_6) \geq 0 . \quad (34)$$

We notice in passing that the fields $C_{1,2}$ appear in the SUSY-breaking scalar potential in the same combination as in the Kähler potential. This particular form may be understood as due to a symmetry under which $C_{1,2} \rightarrow C_{1,2} + i\delta$ in the Kähler potential which is transmitted to the final form of the scalar potential.

It is well known that, for a potential of the generic form (31) (+D-terms), the minimization conditions yield

$$\sin 2\beta = \frac{-2B\mu}{m_{C_1}^2 + m_{C_2}^2 + 2|\mu|^2} . \quad (35)$$

In particular, this relation embodies the boundedness requirement: if the absolute value of the right-hand side becomes bigger than one, this would indicate that the potential becomes unbounded from below. As we have seen, in the class

of models under consideration the particular expressions of the mass parameters lead to the equality (32), which in turns implies $\sin 2\beta = -1$. Thus one finds $\tan \beta = \langle C_2 \rangle / \langle C_1 \rangle = -1$ for any value of $\cos \theta, \Theta_3, \Theta_6$ (and of the other Θ_i 's of course), i.e. for any Goldstino direction.

As an additional comment, it is worth recalling that in previous analyses of the above mechanism for generating μ and B in the String context [4, 5, 23] the value of μ was left as a free parameter since one did not have an explicit expression for the function Z . However, if the explicit orbifold formulae for Z are used, one is able to predict both [7] μ and B reaching the above conclusion⁹.

Now that we have computed explicitly the whole soft terms and the μ parameter, it would be interesting to analyze the dilaton-dominated scenario ($\cos \theta = 0$) because of its predictivity. In particular, from eqs.(8,28,29) we obtain¹⁰

$$\begin{aligned} m_\alpha &= m_{3/2} , \\ M_a &= \pm \sqrt{3} m_{3/2} , \\ A_{\alpha\beta\gamma} &= -M_a , \\ B &= 2m_{3/2} , \\ \mu &= m_{3/2} , \end{aligned} \tag{36}$$

and therefore the whole SUSY spectrum depends only on one parameter $m_{3/2}$. If we would know the particular mechanism which breaks SUSY, then we would be able of computing the superpotential and hence $m_{3/2} = e^K |W|$. Although this is not the case, still this parameter can be fixed from the phenomenological requirement of correct electroweak breaking $2M_W/g_2^2 = \langle H_1 \rangle^2 + \langle H_2 \rangle^2$. Thus at the end of the day we are left with no free parameters. Of course, if in the next future the mechanism which breaks SUSY is known (i.e. $m_{3/2}$ can be explicitly calculated) and the above scenario is the correct one, the value of $m_{3/2}$ should coincide with the one obtained from the phenomenological constraint. In ref.[30] the consistency of the above boundary conditions with the appropriate radiative electroweak symmetry breaking is explored. Unfortunately, it is found that there is no consistency with the measured value of the top-quark mass, namely the mass obtained in this scheme turns out to be too small. A possible way-out to this situation is to assume that also the moduli fields contribute to SUSY breaking since then the soft terms are modified (see eqs.(28,29)) [30]. Of course, this amounts to a departure of the pure dilaton-dominated scenario.

Finally, let us remark that the previous dramatical conclusion in the pure dilaton-dominated limit is also obtained in a different context, namely to avoid low-energy charge and color breaking minima deeper than the standard vacuum [33]. In fact, on these grounds, the dilaton-dominated limit is excluded not only for a μ term generated through the Kähler potential but for any possible mechanism solving the μ problem. The results indicate that the whole free parameter space $(m_{3/2}, B)$ is excluded after imposing the present experimental data on the top mass. The inclusion of a non-vanishing cosmological constant does not improve essentially this situation.

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⁹We should add that situations are conceivable where the above result may be evaded, for example if the physical Higgs doublets are a mixture of the above fields with some other doublets coming from other sectors (e.g. twisted) of the theory.

¹⁰It is worth noticing here that although the value of μ is compactification dependent even in this dilaton-dominated scenario, $\mu = m_{3/2}(\tilde{K}_1 \tilde{K}_2)^{-1/2} Z$ as can be obtained from eq.(23), the result $\mu = m_{3/2}$ will be obtained in any compactification scheme with the following property: $(\tilde{K}_1 \tilde{K}_2)^{-1/2} Z = 1$. Of course, this is the case of orbifolds, where in particular $\tilde{K}_1 = \tilde{K}_2 = Z$, as was shown in eq.(27).

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